The jet power extracted from a magnetized accretion disc

Xinwu Cao

Shanghai Astronomical Observatory, National Astronomical Observatories, Chinese Academy of Sciences,
Shanghai, 200030, China; cxw@center.shao.ac.cn

ABSTRACT
We consider the power of a relativistic jet accelerated by the magnetic field of an accretion disc. It is found that the power extracted from the disc is mainly determined by the field strength and configuration of the field far from the disc. Comparing with the power extracted from a rotating black hole, we find that the jet power extracted from a disc can dominate over that from the rotating black hole. But in some cases, the jet power extracted from a rapidly rotating hole can be more important than that from the disc even if the poloidal field threading the hole is not significantly larger than that threading the inner edge of the disc. The results imply that the radio-loudness of quasars may be governed by its accretion rate which might be regulated by the central black hole mass. It is proposed that the different disc field generation mechanisms might be tested against observations of radio-loud quasars if their black hole masses are available.

Key words: accretion, accretion discs–galaxies: jets–galaxies: nuclei

1 INTRODUCTION

The current most favoured models of powering active galactic nuclei (AGNs) involve gas accretion onto a massive black hole, though the details are still unclear. Relativistic jets have been observed in many radio-loud AGNs and are believed to be formed very close to the black holes. The spin energy of a black hole might be extracted to fuel the jet by magnetic fields supported by a surrounding accretion disc (Blandford & Znajek 1977; Blandford-Znajek process). This process has widely been believed to be the major mechanism that powers radio jets in AGNs (Begelman, Blandford & Rees 1984; Rees et al. 1982; Wilson & Colbert 1995; Moderski & Sikora 1996). However, Ghosh & Abramowicz (1997) doubted the importance of the Blandford-Znajek process. For a black hole of a given mass and angular momentum, the strength of the Blandford-Znajek process depends crucially on the strength of the poloidal field threading the horizon of the hole. They argued that the strength of the field threading a black hole has been overestimated. The magnetic field threading a hole should be maintained by the currents situated in the inner region of a surrounding accretion disc. Livio, Ogilvie & Pringle (1999; hereafter L99) re-investigated the problem and pointed out that even the calculations of Ghosh & Abramowicz (1997) have overestimated the power of the Blandford-Znajek process. Ghosh & Abramowicz (1997) have overestimated the strength of the large-scale field threading the inner region of an accretion disc, and then the power of the Blandford-Znajek process. They argued that the jet power extracted from an accretion disc dominates over the power extracted by the Blandford-Znajek process (L99).

L99 estimated the maximal electromagnetic output from an accretion disc. In this case, the toroidal field component is of the same order of the poloidal field component at the disc surface due to the instability of a predominantly toroidal field (Biskamp 1993; Livio & Pringle 1992). Apart from the strength of the field threading a disc, the acceleration of the jet is also governed by the magnetic field configuration and the structure of the disc (Shu 1991; Ogilivie & Livio 1998, 2000; Cao & Spruit 2001; for relativistic jets near rotating black holes, see Koide, Shibata & Kudoh 1999; Koide et al. 2000). It is unclear how much power can actually be extracted from a magnetized accretion disc without considering its field configuration. In this work, we extend L99's work to explore some factors that affect the power of a relativistic jet extracted from an accretion disc.

2 THE POWER OF A JET EXTRACTED FROM AN ACCRETION DISC

For a relativistic jet accelerated by the magnetic field of the disc, the Alfvén velocity is (Michel 1969; Camenzind 1986)

$$v_{\rm A} = \frac{B_{\rm p}^{\rm A}}{(4\pi\rho_{\rm A}\gamma_{\rm j})^{1/2}}, \tag{1}$$

where $B_{\rm p}^{\rm A}$ and $\rho_{\rm A}$ are the poloidal field strength and density of the jet at Alfvén point, and $\gamma_{\rm j}$ is the Lorentz factor of the bulk motion of the jet. Equation (1) is a good approximation, since the Alfvén point is usually not very close to the rotating black hole (Camenzind 1986). The value of the Alfvén velocity is of the order $v_{\rm A} \sim R_{\rm A}\Omega(R_{\rm d})$, where $R_{\rm d}$ is the radius of the field footpoint at the disc surface, and $R_{\rm A}$ is the Alfvén radius of the jet along the field line.

Mass and magnetic flux conservation along a magnetic field line requires

$$\frac{\dot{m}_{\rm jet}}{B_{\rm pd}} \simeq \frac{\rho_{\rm A} v_{\rm A}}{B_{\rm p}^{\rm A}}, \tag{2}$$

where $\dot{m}_{\rm jet}$ is the mass loss rate in the jet from unit surface area of the disc, and $B_{\rm pd}$ is the poloidal field strength at the disc surface.

Since the Alfvén radius $R_{\rm A}$ depends on the global field configuration, not only the strength of the field at the disc surface, we assume the magnetic field far from the disc surface to be roughly self-similar:

$$B_{\rm p}^{\rm A} \sim B_{\rm pd} \left(\frac{R_{\rm A}}{R_{\rm d}} \right)^{-\alpha}, \tag{3}$$

where the self-similar index $\alpha > 1$ (Blandford & Payne 1982; Lubow, Papaloizou & Pringle 1994). A self-similar index $\alpha = 4$ is adopted in the calculations of Lubow et al. (1994).

The Lorentz factor of the jet is

$$\gamma_{\rm j} \simeq \left[1 - \left(\frac{v_{\rm A}}{c} \right)^2 \right]^{-\frac{1}{2}}. \tag{4}$$

Combining Eqs (1)-(4), the mass loss rate in the jet from unit surface area of the disc is

$$\dot{m}_{\rm jet} = \frac{B_{\rm pd}^2}{4\pi c} \left(\frac{R_{\rm d}\Omega(R_{\rm d})}{c} \right)^\alpha \frac{\gamma_{\rm j}^\alpha}{(\gamma_{\rm j}^2 - 1)^{\frac{1-\alpha}{2}}}. \tag{5}$$

The power of the jet accelerated by the magnetized disc is

$$L_{\rm d} = 2\pi \int (\gamma_{\rm j} - 1)\dot{m}_{\rm jet}c^2 R_{\rm d} dR_{\rm d}$$

$$= 2\pi c \int \frac{B_{\rm pd}^2}{4\pi} \left(\frac{R_{\rm d}\Omega(R_{\rm d})}{c} \right)^\alpha f(\alpha, \, \gamma_{\rm j}) R_{\rm d} dR_{\rm d}, \tag{6}$$

where the function $f(\alpha, \, \gamma_{\rm j})$ is given by

$$f(\alpha, \, \gamma_{\rm j}) = \frac{(\gamma_{\rm j} - 1)\gamma_{\rm j}^\alpha}{(\gamma_{\rm j}^2 - 1)^{\frac{1-\alpha}{2}}}. \tag{7}$$

For a given $\gamma_{\rm j}$, Eq. (6) reduces to

$$L_{\rm d} = 2\pi c f(\alpha, \, \gamma_{\rm j}) \int \frac{B_{\rm pd}^2}{4\pi} \left(\frac{R_{\rm d}\Omega(R_{\rm d})}{c} \right) R_{\rm d} dR_{\rm d}, \tag{8}$$

in the case of $\alpha = 1$. The power of the jet $L_{\rm d}$ decreases with the self-similar index α, since the term $R_{\rm d}\Omega(R_{\rm d})/c$ in Eq. (6) is always less than unity. Thus, Eq. (8) gives an upper limit on the power of a relativistic jet for given field strength at the disc surface and Lorentz factor $\gamma_{\rm j}$ of the jet ($\gamma_{\rm j} \gg 1$; see Fig. 1). Note that Eq. (8) is almost same as Eq. (6) in L99 for the maximal electromagnetic output from the disc, except an additional term of $f(\alpha, \, \gamma_{\rm j})$ appears here. Here we have obtained a similar conclusion as that in L99, but in a rather different way.

The toroidal field strength at the disc surface is

$$B_{\phi \rm d} = \frac{4\pi}{B_{\rm pd}} \frac{\dot{m}_{\rm jet}c^2}{R_{\rm d}\Omega(R_{\rm d})} (\gamma_{\rm j} - 1). \tag{9}$$

Substituting Eq. (5) into Eq. (9), we have

$$\frac{B_{\phi \rm d}}{B_{\rm pd}} = \left(\frac{R_{\rm d}\Omega(R_{\rm d})}{c} \right)^{\alpha - 1} f(\alpha, \, \gamma_{\rm j}). \tag{10}$$

It is obvious that the ratio $B_{\phi \rm d}/B_{\rm pd}$ at the given radius $R_{\rm d}$ is determined by the self-similar index α and the Lorentz factor $\gamma_{\rm j}$, i.e., the ratio only depends on the jet properties, not on the disc properties.

2.1 The magnetic field strength at the disc surface

In order to complete the estimate on the power of the jet extracted from an accretion disc, we need to know the field strength at the disc surface. The origin of the ordered field that is assumed to thread the disc is still unclear. The strength of the field at the disc surface is usually assumed to scale with the pressure of the disc, as done in Ghosh & Abramowicz (1997). However, L99 have estimated the strength of the large-scale field threading the disc based on the results of some numerical simulations on dynamo mechanisms in accretion discs (Tout & Pringle 1996; Amitage 1998; Romanova et al. 1998). In this work, we mainly follow L99's estimate on the field strength at the disc surface. The physical implications of these two different estimates on the field strength at the disc surface will be discussed in Sects. 4 and 5.

As L99, the strength of the magnetic field produced by dynamo processes in the disc is given by

$$\frac{B_{\rm dynamo}^2}{4\pi} \sim \frac{W}{2H}, \tag{11}$$

where W is the integrated shear stress of the disc, and H is the scale-height of the disc. For a relativistic accretion disc, the integrated shear stress is given by Eq. (5.6.14a) in Novikov & Thorne (1973, hereafter NT73). Equation (9) can be re-written as

$$B_{\rm dynamo} = 3.56 \times 10^8 r_{\rm d}^{-3/4} m^{-1/2} A^{-1} B E^{1/2} {\rm gauss}, \tag{12}$$

where the dimensionless quantities are defined by

$$r_{\rm d} = \frac{R_{\rm d}}{R_{\rm G}}, \quad R_{\rm G} = \frac{GM}{c^2}, \quad m = \frac{M}{M_\odot}, \tag{13}$$

and A, B and E are general relativistic correction factors defined in NT73.

In standard accretion disc models, the angular velocity of the matter in the disc is usually very close to Keplerian velocity. For a relativistic accretion disc surrounding a rotating black hole, the Keplerian angular velocity is given by

$$\Omega(r_{\rm d}) = \frac{M^{1/2}}{r_{\rm d}^{3/2} + a}, \tag{14}$$

where a is dimensionless specific angular momentum of a rotating black hole.

The large-scale field can be produced from the small-scale field created by dynamo processes as $B(\lambda) \propto \lambda^{-1}$ for the idealized case, where λ is the length scale of the field (Tout & Pringle 1996; Romanova et al. 1998). The large-scale field threading the disc is related with the field produced by dynamo processes approximately by (L99)

$$B_{\rm pd} \sim \left(\frac{H}{R_{\rm d}} \right)_{\rm max} B_{\rm dynamo}. \tag{15}$$

The dimensionless scale-height of a disc $h = H/R_{\rm d}$ is in principle a function of $R_{\rm d}$, and it reaches a maximal value in the inner region of the disc (Laor & Netzer 1989). We adopt the maximal value of $H/R_{\rm d}$ here.

The scale-height of the disc $H/R_{\rm d}$ is given by (Laor & Netzer 1989)

$$\frac{H}{R_{\rm d}} = 26.345 \dot{m}_{\rm acc} r_{\rm d}^{-1} c_2, \tag{16}$$

where the coefficient c_2 is defined in NT73. The dimensionless accretion rate is defined by (Laor & Netzer 1989)

$$\dot{m}_{\rm acc} = \frac{\dot{M}_{\rm acc}}{1.626 \times 10^{15} m \,\, {\rm g \,\, s}^{-1}}. \tag{17}$$

We use Eqs. (12) and (15), the power of the jet accelerated from a magnetized disc is available by integrating Eq. (6), if some parameters: m, $\dot{m}_{\rm acc}$, a, α and $\gamma_{\rm j}$ are specified.

3 THE RELATIVE IMPORTANCE OF THE BLANDFORD-ZNAJEK PROCESS

As discussed in L99, the power extracted from a rotating black hole by the Blandford-Znajek process is determined by the hole mass M, the spin of the hole a, and the strength of the poloidal field threading the horizon of a rotating hole $B_{\rm ph}$:

$$L_{\rm BZ}({\rm max}) \sim \frac{B_{\rm ph}^2}{4\pi} \pi R_{\rm h}^2 a^2 c, \tag{18}$$

where $R_{\rm h}$ is the radius of the black hole horizon. The strength of the poloidal field threading the hole is a crucial factor determining the power extracted from the hole. As argued in Ghosh & Abramowicz (1997) and L99, the field threading the black hole is generated by the currents situated in the disc rather than in the hole. Thus, it is realized that the field threading the hole should not be significantly larger than the field threading the inner edge of the disc Here we use a parameter ζ to relate $B_{\rm ph}$ with $B_{\rm pd}(r_{in})$:

$$B_{\rm ph} = \zeta B_{\rm pd}(r_{\rm in}), \tag{19}$$

where $\zeta \geq 1$, $r_{\rm in}$ is the radius of the inner edge of the disc. For a standard accretion disc, $r_{\rm in} = r_{\rm ms}$, where $r_{\rm ms}$ is the radius of the minimal stable circular orbit.

The Numerical simulations on the Blandford-Znajek process show that it can operate near its maximum power output (Komissarov 2001). Using Eqs. (6), (18) and (19), we can compare the jet power extracted from a magnetized disc with that by the Blandford-Znajek process. It is worth noting that the ratio of the jet power extracted from the disc to that from a rotating black hole is only determined by the parameters α, $\gamma_{\rm j}$ and ζ. The relative importance of these two processes is independent of the mechanism for producing the disc field.

4 RESULTS

We plot the function $f(\alpha, \gamma_{\rm j})$ varying with $\gamma_{\rm j}$ for different values of α in Fig. 1. The final velocity of a magnetically driven jet depends on the strength and global configuration of the field threading the disc, and the structure of the disc that provides boundary conditions for the jet (Shu 1991; Ogilivie & Livio 1998, 2000; Cao & Spruit 1994, 2001). In this work, we focus on relativistic jets which are widely observed in AGNs. To avoid being involved in the complexity of the jet acceleration problem, we use the Lorentz factor $\gamma_{\rm j}$ of a magnetically accelerated jet as a parameter to describe the problem. Now the problem is described by the parameters m, a, $\dot{m}$, α and $\gamma_{\rm j}$.

The mass loss rate in the jet from unit surface area of the disc is plotted in Figs. 2 and 3 for the non-rotating and rapidly rotating black holes respectively. We plot the total mass loss rate in the jet and the power of the jet as functions of $\gamma_{\rm j}$ for different values of α in Figs. 4 and 5 respectively. We calculate the ratio of the toroidal field strength to the poloidal one at the disc surface using Eq. (10), and the results are present in Fig. 6.

Using Eq. (18), we can estimate the maximal power output from a rotating black hole by the Blandford-Znajek process. The power outputs as functions of a are present in Fig. 7 for different
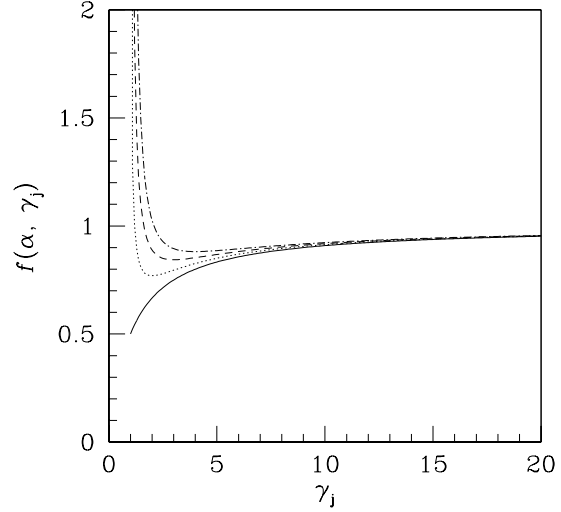

Figure 1. The function $f(\alpha, \gamma_{\rm j})$ varies with $\gamma_{\rm j}$ for different values of α: $\alpha = 1$(solid line), 2(dotted), 3(dashed) and 4(dash-dotted).

values of α and ζ. We plot the power outputs varying with accretion rate $\dot{m}_{\rm acc}$ based on L99's estimate of the disc field strength in Fig. 8. Finally, using Eqs. (6), (18) and (19), we plot the relative importance of these two power extract processes in the parameter space (α-a) in Fig. 9.

4.1 The ratio of jet power to accretion luminosity

The power of the jet extracted from a disc or a rotating black hole depends mainly on the field strength at the disc surface, since the field threading the hole is generated by the currents situated in the disc. The ratio of jet power to accretion luminosity is therefore related with the mechanism for generating the disc field.

In the case of the large-scale field being produced from the small-scale field created by dynamo processes, the large-scale field strength is estimated by Eq. (15). As we know, $H/R_{\rm d} \propto \dot{m}_{\rm acc}$ and $L_{\rm acc} \propto \dot{m}_{\rm acc} m$, then we obtain a linear relation between $\dot{m}_{\rm acc}$ and $L_{\rm jet}/L_{\rm acc}$ by using Eqs. (6) (see Fig. 8):

$$L_{\rm jet}/L_{\rm acc} \propto \dot{m}_{\rm acc}, \tag{20}$$

which is independent of the black hole mass. Note that the strength of the field threading the hole is related with that threading the disc (see Eqs. (18) and (19)), the relation (20) is also valid for jet power extracted from the rotating hole.

The similar analysis can be performed for the case that the magnetic field strength scales with the pressure of the disc. The strength of the field at the disc surface is

$$B_{\rm pd}(r_{\rm d}) \propto r_{\rm d}^{-3/4} m^{-1/2} A^{-1} B E^{1/2}, \tag{21}$$

which is in dependent of the accretion rate (NT73).

From Eqs. (6), (18) and (21), we obtain a relation (Ghosh & Abramowicz 1997)

$$L_{\rm jet}/L_{\rm acc} \propto \dot{m}_{\rm acc}^{-1}, \tag{22}$$

which holds for a jet accelerated either from a disc or a rotating hole. It does not depend on the black hole mass either.

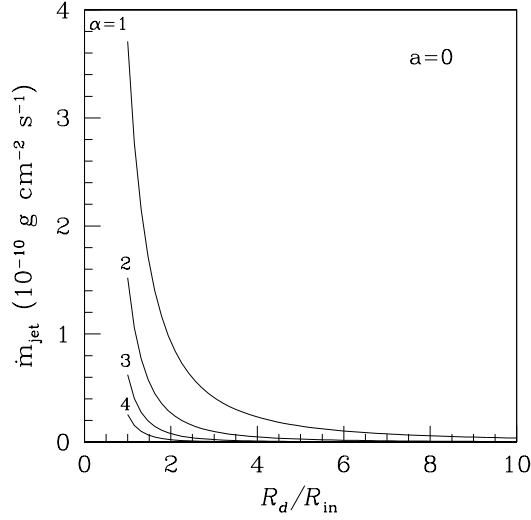**Figure 2.** The mass loss rate in the jet from unit surface area of the disc surrounding a non-rotating black hole for different values of α. The hole mass $m = 10^9$, $\gamma_j = 10$ and $\dot{m}_{acc} = 0.05$ are adopted.
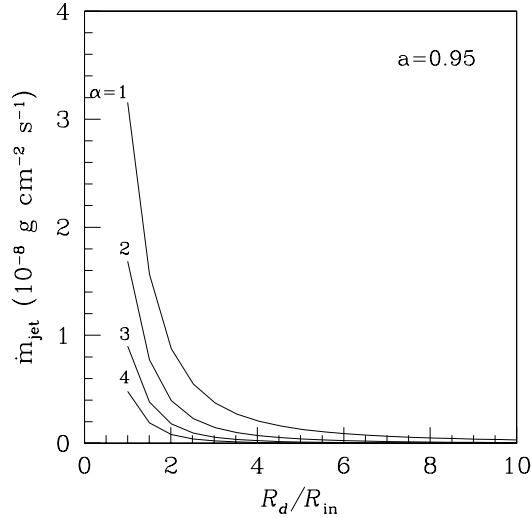

Figure 4. The total mass loss rate in the jet as functions of γ_j for different values of α for the field. The black hole spin parameter $a = 0.95$ and $\dot{m}_{acc} = 0.05$ are adopted.

Figure 3. Same as Fig. 2, but for a rapidly rotating black hole: $a = 0.95$.

Figure 5. The power of the jet accelerated by a magnetized disc as functions of γ_j for different values of α. The left panel is for a non-rotating black hole, while the right panel is for a rapidly rotating black hole: $a = 0.95$. The accretion rate $\dot{m}_{acc} = 0.05$ is adopted.

5   DISCUSSION OF THE RESULTS

Figure 1 shows that the function $f(\alpha, \gamma_j)$ varies slowly with γ_j except for $\gamma \to 1$. The function $f(\alpha, \gamma_j)$ is almost a constant while $\gamma_j > 5$. From Eq. (6), it is then found that the power of a relativistic jet accelerated by a magnetized accretion disc is insensitive to its Lorentz factor γ_j. This is confirmed by the results plotted in Fig. 5. The jet power varies with γ_j in a range less than 10 % for $\gamma_j > 5$. It implies that relativistic jets with different Lorentz factors γ_j accelerated by magnetized discs have almost same jet power, if all other physical parameters are fixed. We can therefore adopt a typical value of γ_j in our calculations, which will not change the main results on jet power, though the terminal velocity of the magnetically driven jet could be given in principle only after the structure of the disc and its field strength and configuration are supplied. This
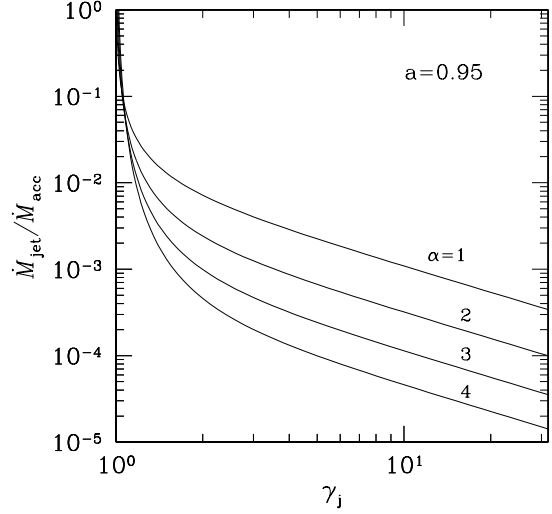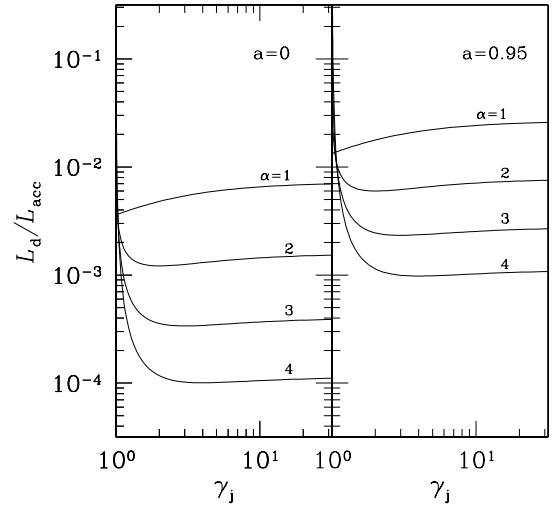
has simplified our present investigation, since we are interested in the power of relativistic jets extracted from discs.

The power of a jet decreases with the self-similar index α (see Fig. 5). A large α implies that the poloidal field strength decays rapidly with radius along a field line. From Fig. 4, we see that less mass is loaded into the jet in the case of a larger α, and the power extracted from the disc is then reduced. The power of the jet is sensitive to the field configuration far from the disc surface.

We find that the mass loss rate in the jet from unit surface area of the disc decreases with radius (see Figs. 2 and 3). Almost all mass carried by the jet is from the region of the disc: $R_d < 5\, R_{in}$, i.e., jets are mainly accelerated from the inner region of the disc
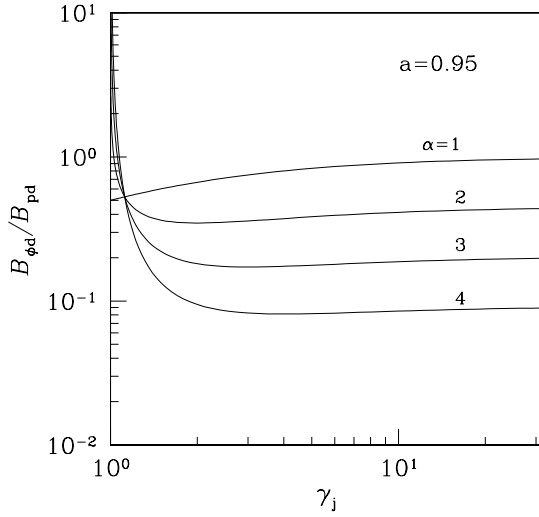
Figure 6. The ratio of the toroidal field strength to the poloidal one at the disc surface as functions of $\gamma_{\rm j}$ for different values of α at $R_{\rm d} = 2R_{\rm in}$. The hole spin parameter $a = 0.95$ is adopted.

Figure 8. The power outputs as functions of $\dot{m}_{\rm acc}$ for different values of α and ζ. The solid lines represent the power output from an accretion disc, while the dotted lines represent the power output of the Blandford-Znajek process. The black hole spin parameter $a = 0.95$ is adopted.

Figure 7. The power outputs as functions of a for different values of α and ζ. The solid lines represent the power output from an accretion disc, while the dotted lines represent the power output of the Blandford-Znajek process. The accretion rate $\dot{m}_{\rm acc} = 0.05$ is adopted.

Figure 9. The relative importance of two power extract processes in parameter space (α-a) for different values of ζ: $\zeta = 1$(solid line), 2(dotted) and 3(dashed). $\gamma_{\rm j} = 10$ is adopted.

where most gravitational energy of accretion matter is released. The total mass loss rate in a relativistic jet can be neglected compared with the accretion rate of the disc (usually $\dot{M}_{\rm jet} < 10^{-2}\dot{M}_{\rm acc}$, see Fig. 4). A jet with high mass loss rate results in a low terminal velocity of the jet and wound-up field lines (Figs. 4 and 6), which is consistent with the results of Cao & Spruit (1994). As pointed out in Sect. 2, our calculation of the jet power extracted from a disc can reproduce L99's result while $\alpha = 1$ is adopted (see Eq. (8)). In L99's calculation, $B_{\phi \rm d} = B_{\rm pd}$ at the disc surface is assumed, so it is not surprised to find that $B_{\phi \rm d} \simeq B_{\rm pd}$ in the case of $\alpha = 1$ in our calculations (see Fig. 6). For a large α, the ratio $B_{\phi \rm d}/B_{\rm pd}$ is low (compare the results in Fig. 6 for different values of α).
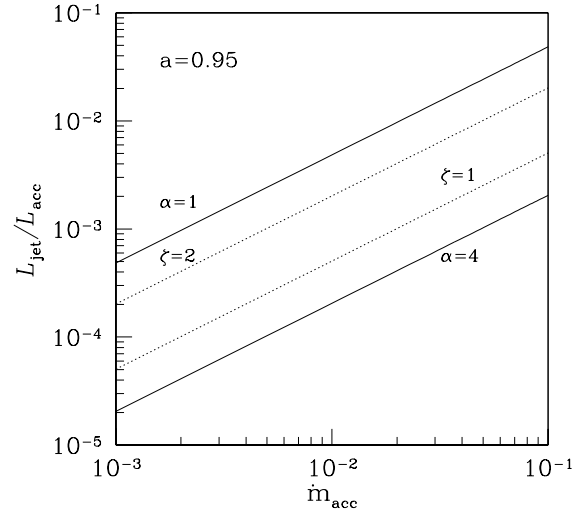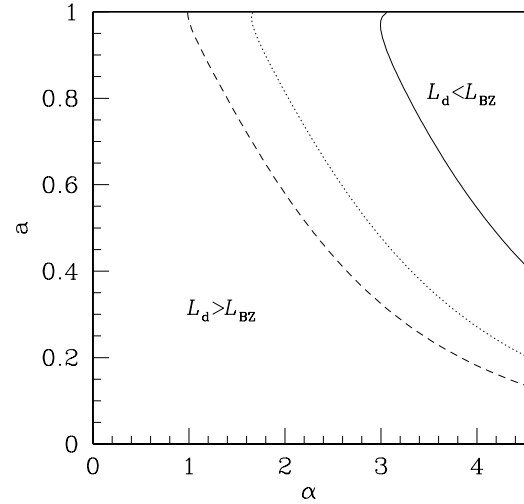
In our present calculations, the field configuration far from the disc is described by the self-similar index α. In the case of $\alpha = 1$, it is found in Fig. 7 that the power extracted from the disc always dominates over the maximal power of the Blandford-Znajek process, if $\zeta < 3$, i.e., the strength of the poloidal field threading the hole is less than three times of that threading the inner edge of the disc. If $\zeta = 1$ is assumed, i.e., $B_{\rm ph} = B_{\rm pd}(R_{\rm in})$, $L_{\rm d}$ is an order of magnitude larger than $L_{\rm BZ}$ even for an extreme Kerr black hole. Thus the conclusion given by L99 that the maximal power extracted from a disc dominates over that from a rapidly rotating black hole is confirmed, while our calculations are not limited to the extreme case $\alpha = 1$. We present the results in Fig. 7 to compare the power

outputs by these two processes for different values of α and ζ. It is found that the Blandford-Znajek process would be important for a large α or/and ζ. From Eqs. (6) and (18), we find that this conclusion is independent of the mechanism for producing disc fields. The conclusion is also independent of the hole mass m and accretion rate $\dot{m}_{\rm acc}$, since the hole field strength is only governed by the disc field strength.

If the large-scale field can be produced from the small-scale field created by dynamo processes as done by L99, we find that linear relations are present: $L_{\rm jet}/L_{\rm acc} \propto \dot{m}_{\rm acc}$, for a jet accelerated either from a disc or a rotating hole. (see Fig. 8 and Eq. (20)). This relation is quite different from the situation for the field that is assumed to scale with the pressure in the disc, where $L_{\rm jet}/L_{\rm acc} \propto \dot{m}_{\rm acc}^{-1}$ is present.

The ratio $L_{\rm jet}/L_{\rm acc}$ can be related with an observational quantity of quasars: radio-loudness R, which is defined as $R \equiv f_\nu(5{\rm GHz})/f_\nu(4400\text{Å})$. The fact that the ratio $L_{\rm jet}/L_{\rm acc}$ is independent of the black hole mass may imply that the radio-loudness R is not related directly with the black hole mass. The radio-loudness R may probably governed by the accretion rate. The accretion rate is then regulated by the central black hole mass in some way (Yi 1996; Salucci et al. 1999; Haiman & Menou 2000; Kauffmann & Haehnelt 2000). The relations found between radio-loudness and the central black hole masses in quasars (McLure et al. 1999; Laor 2000; Gu, Cao & Jiang 2001; Ho 2002) may reflect intrinsic relations between radio-loudness and accretion rate which is regulated by the black hole mass (Boroson 2002). The different relations between $L_{\rm jet}/L_{\rm acc}$ and $\dot{m}_{\rm acc}$ are expected for different mechanisms of producing the disc field. The different disc field generation mechanisms can therefore be tested against observations of radio-loud quasars if their black hole masses are available. However, whether the jet power is extracted from a disc or a rotating black hole cannot be tested in this way, since the calculations predict the same relation between $L_{\rm jet}/L_{\rm acc}$ and $\dot{m}_{\rm acc}$ for these two jet acceleration mechanisms.

In Fig. 9, we compare the relative importance of these two power extract processes in the parameter space ($\alpha - a$). Only for those cases the strength of the poloidal field of the disc decays slowly with radius along the field line (small α), the power extracted from the disc dominates over that extracted by the Blandford-Znajek process even for a rapidly rotating hole. In the right-upper corner of the figure, i.e., the case with a large α and a, the power extracted from the disc is greater than the power of the Blandford-Znajek process. Otherwise, the power of the Blandford-Znajek process can be neglected compared with that from the disc.

The acceleration of the jet is governed by the structure of the disc and the strength and configuration of the field threading the disc. The physical properties of the jet, such as the power and terminal velocity of the jet, can be determined if all these physical factors of the disc and its field are known. Instead of solving a set of MHD equations describing the disc-jet system, a self-similar index α is employed to describe the field far above the disc surface in present work. The physics of how the jet acceleration is governed by the detailed structure of the disc near a rapidly rotating black hole has not been included in this work. It would be the main limitation of the present work.

ACKNOWLEDGMENTS

REFERENCES

Amitage P. J., 1998, ApJ, 501, L189
Begelman M. C., Blandford R. D., Rees M. J., 1984, Rev. Mod. Phys., 56, 255
Biskamp D., 1993, Nonlinear magnetohydrodynamics, Cambridge Univ. Press, Cambridge
Blandford R. D., Payne D. G., 1982, MNRAS, 199, 883
Blandford R. D., Znajek R. L., 1977, MNRAS, 179, 433
Boroson T. A., 2002, ApJ, 565, 78
Camenzind M., 1986, A&A, 156, 137
Cao X., Spruit H. C., 1994, A&A, 287, 80
Cao X., Spruit H. C., 2002, A&A, in press
Ghosh P., Abramowicz M. A., 1997, MNRAS, 292, 887
Gu M., Cao X., Jiang D. R., 2001, MNRAS, 327, 1111
Haiman Z., Menou K., 2000, ApJ, 531, 42
Ho L. C., 2002, ApJ, 564, 120
Kauffmann G., Haehnelt M., 2000, MNRAS, 311, 576
Koide S., Shibata K., Kudoh T., 1999, ApJ, 522, 727
Koide S., Meier D., Shibata K., Kudoh T., 2000, ApJ, 536, 668
Komissarov S. S., 2001, MNRAS, 326, L41
Laor A., 2000, ApJ, 543, L111
Laor A., Netzer H., 1989, MNRAS, 238, 897
Livio M., Ogilivie G. I., Pringle J. E., 1999, ApJ, 512, 100(L99)
Livio M., Pringle J. E., 1992, MNRAS, 259, 23
Lubow S. H., Papaloizou J. C. B., Pringle J. E., 1994, MNRAS, 268, 1010
McLure R.J., Kukula M.J., Dunlop J.S., Baum S.A., O'Dea C.P., Hughes D.H., 1999, MNRAS, 308, 377
Michel F. C., 1969, ApJ 158, 727
Moderski R., Sikora M., 1996, A&AS 120, 591
Novikov I., Throne K. S., 1973, in Black holes, eds de Witt C. and de Witt B., Gordon & Breach, New York(NT73)
Ogilvie G.I., Livio M., 1998, ApJ, 499, 329
Ogilvie G.I., Livio M., 2001, ApJ, 553, 158
Rees M. J., Begelman M. C., Blandford R. D., Phinney E. S., 1982, Nat, 295, 17
Romanova M. M., Ustyugova G. V., Koldoba A. V., Chechetkin V. M., Lovelace R. V. E., 1998, ApJ, 500, 703
Salucci P., Szuszkiewicz E., Monaco P., Danese L., 1999, MNRAS, 307, 637
Shu F.H., 1991, in Lambert D.L., ed, Frontiers of stellar evolution, Astronomical Society of the Pacific Conference Series (ASP), p.23
Tout C. A., Pringle J. E., 1996, MNRAS, 281, 219
Wilson A. S., Colbert E. J. M., 1995, ApJ, 438, 62
Yi I., 1996, ApJ, 473, 645